Flavor Structure of Scherk-Schwarz Supersymmetry Breaking and Constraints from Low Energy Processes [a]

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ABSTRACT

5D orbifold has two attractive features: quasi-localized matter fields can naturally reproduce hierarchical Yukawa coupling, while supersymmetry breaking is inherently built in by the Scherk-Schwarz mechanism. We examine the consequence of this quasi-localization and the Scherk-Schwarz supersymmetry breaking in low energy flavor violating processes, under the assumption that physics below the compactification scale is described by the minimal supersymmetric standard model (MSSM) and find $BR(\mu \to e, \gamma)$ and ϵ_K impose stringent constraint on flavor structure of the Scherk-Schwarz supersymmetry breaking. Chirality measurement in lepton flavor violating processes is crucial to deduce the surviving flavor structure.

Localization of matter fields along the extra-dimension can naturally reproduce observed Yukawa hierarchy [1], while the Scherk-Schwarz mechanism provides a simple alternative way to break supersymmetry [2]. 5D orbifold S^1/Z_2 is a minimal realization of both of these attractive features of introducing new dimensions. Let us first define the model using 5D bulk action on S^1 in terms of 4D superfield representation [3],

$$\begin{aligned} S_{bulk} \;&=\; \int d^4x \int dy \left[\int d^4\theta \left\{ Re(T)\left(\Phi_I \Phi_I^* + \Phi_I^c \Phi_I^{c*}\right) + \frac{1}{g_{5a}^2}\frac{1}{Re(T)}\left(\partial_5 V - \sqrt{2}Re(\chi)\right)^2 \right\} \right. \\ &\quad \left. + \left(\int d^2\theta\, \Phi_I^c \left(\partial_y - \frac{1}{\sqrt{2}}\chi + M_I T\right)\Phi_I + \frac{1}{4g_{5a}^2} T W^{a\alpha} W_\alpha^a + \text{h.c.} \right) \right], \end{aligned} \tag{1}$$

where 4D chiral multiplets (Φ_I, Φ_I^c) and 4D vector and chiral multiplets (V, χ) constitute 5D hyper and vector multiplets respectively. 5D supersymmetry forbids bulk Yukawa coupling. Here, radius modulus R of S^1, satisfying $x^5 = Ry$ $(0 \le y < 2\pi)$, is promoted to the radion chiral multiplet T so that 4D N=1 supersymmetry is manifest. In 5D supergravity, it is identified as a part of 5D supergravity multiplet [2]. Because we are interested in 4D chiral theories, we orbifold S_1 by Z_2 imposing boundary conditions,

$$V^a(-y) = V^a(y), \quad \chi^a(-y) = -\chi^a(y), \quad \Phi_I(-y) = \Phi_I(y), \quad \Phi_I^c(-y) = -\Phi_I^c(y), \tag{2}$$

so that 4D $N = 2$ supersymmetry is broken to $N = 1$. After orbifolding, the surviving bulk mass should be regarded as a kink mass, $M_I \to M_I \epsilon(y)$ respecting Z_2. Then the zero mode equation for Φ_I fermion component, $\psi_I = \chi(x)\tilde{\phi}_{0I}(y)$ is given by $(\,\partial_y + M_I T \epsilon(y)\,)\tilde{\phi}_{0I} = 0$, yielding the zero mode wavefunction $\tilde{\phi}_{0I} \propto e^{-M_I T|y|}$, which shows that it is quasi-localized

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at $y=0$ if $M_I>0$, and at $y=\pi$ if $M_I<0$. Inspired by this observation, we assume Higgs fields exist on the fixed point at $y=0$ and introduce the following brane action,

$$S_{brane}=\int d^4x\int d^2\theta\, H\int dy\,\delta(y)\lambda^{(5)}_{IJ}\Phi_I\Phi_J+\text{h.c}\,, \tag{3}$$

After redefinition of the chiral fields, $\Phi_I\;\rightarrow\;e^{M_IT|y|}\Phi_I,\;\Phi^c_I\;\rightarrow\;e^{-M_IT|y|}\Phi^c_I$, and appropriate renormalization, we can extract the 4D effective action for Φ_I zero mode,

$$S_{4D}=\int d^4x\left[\int d^4\theta\,(\Phi_I\Phi^*_I+HH^*)+\left(\int d^2\theta\, y_{IJ}H\Phi_I\Phi_J+\text{H.c.}\right)\right], \tag{4}$$

where we use the same symbol Φ_I for the corresponding canonical zero-modes and,

$$y_{IJ}=\lambda^{(4)}_{IJ}\sqrt{\frac{N_IN_J}{(\epsilon^{-2N_I}-1)(\epsilon^{-2N_J}-1)}},\quad\text{for}\quad\lambda^{(4)}_{IJ}=\frac{\ln(1/\epsilon)}{\pi R}\lambda^{(5)}_{IJ},\quad N_I=-\frac{\pi R}{\ln(1/\epsilon)}M_I. \tag{5}$$

Here, we choose $\epsilon\;\equiv\;$ Cabibbo angle $\;\approx\;0.2$ for later convenience. Naive dimensional argument shows $\lambda^{(4)}_{IJ}\sim\mathcal{O}(1)$ while hierarchical y_{IJ} is naturally realized for $N_I>0$.

So far the 4D effective theory is supersymmetric, however, the remaining supersymmetry can be generally broken by non-vanishing radion F-term, F^T [4]. The resultant soft terms are easily obtained by standard calculus as follows,

$$S_{soft}=-\int d^4x\,\Big(\frac{1}{2}m^2_{IJ}\phi^*_I\phi_J+A_{IJ}\phi_I\phi_Jh+\frac{1}{2}M_a\lambda^a\lambda^a+\text{h.c.}\Big), \tag{6}$$

$$\begin{aligned}&M_a=M_{1/2}=-\frac{F^T}{2R}\,,\qquad A_{IJ}=2y_{IJ}\ln(1/\epsilon)\frac{F^T}{2R}\left(\frac{N_I}{1-\epsilon^{2N_I}}+\frac{N_J}{1-\epsilon^{2N_J}}\right),\\ &m^2_{I\bar{J}}=\delta_{IJ}\left(2\ln(1/\epsilon)\frac{N_I}{\epsilon^{N_I}-\epsilon^{-N_I}}\frac{F^T}{2R}\right)^2,\end{aligned} \tag{7}$$

where lower-case fields denote scalar components of the chiral multiplets and λ^a represents gaugino. In $F^T/R<<1$ limit, this mass spectrum exactly matches with that of the Scherk-Schwarz supersymmetry breaking (SS SUSY breaking), where twisted boundary condition is imposed on $SU(2)_R$ over S^1/Z_2 [5]. It is highly dependent on the flavor structure given in (5), therefore becomes a source of dangerous flavor changing processes at the electro-weak scale. In the following, we present a brief summary of systematic analysis of this low energy constraint on the flavor structure of the SS SUSY breaking.

The above Yukawa couplings are quite similar to the Yukawa couplings in Frogatt-Nielsen models with N_I being identified as the $U(1)_F$ charges. More explicitly, assuming that mass spectrum below the compactification scale is given by that of the minimal supersymmetric standard model (MSSM),

$$y^u_{ij}\;\simeq\;\lambda^{u\,(4)}_{ij}\epsilon^{X^u_i+X^q_j}\,,\quad y^d_{ij}\;\simeq\;\lambda^{d(4)}_{ij}\epsilon^{X^d_i+X^q_j}\,,\quad y^\ell_{ij}\;\simeq\;\lambda^{\ell(4)}_{ij}\epsilon^{X^e_i+X^\ell_j}\,, \tag{8}$$

where u_i, d_i (e_i) denote the $SU(2)_L$ singlet quarks (leptons) and q_i (ℓ_i) represents the doublet quarks (leptons). The effective flavor charge X_I is given by $X_I=N_I$ ($X_I=0$)

for $N_I \geq 0$ ($N_I \leq 0$). The above soft terms also can be well approximated by X_I and N_I,

$$A^u_{ij} \simeq M_0(X^u_i + X^q_j)\, y^u_{ij}, \quad A^d_{ij} \simeq M_0(X^d_i + X^q_j)\, y^d_{ij}, \quad A^\ell_{ij} \simeq M_0(X^e_i + X^\ell_j)\, y^\ell_{ij},$$

$$m^{2(\tilde{\psi})}_{i\bar{j}} \simeq \delta_{i\bar{j}} M_0^2 \begin{cases} |N^\psi_i|^2 \epsilon^{2|N^\psi_i|} & (\, N^\psi_i \neq 0\,) \\ 1/[2\ln(1/\epsilon)]^2 & (\, N^\psi_i = 0\,) \end{cases} \quad \text{for} \quad M_0 = -2M_{1/2}\ln(1/\epsilon). \tag{9}$$

To discuss the low energy observables, it is convenient to move from the above definition basis to the SCKM basis where quarks and leptons have diagonal mass matrices, *e.g.* $y^\ell_{ij} \longrightarrow (\overline{V}_e)_{ik} y^\ell_{k\ell} (\overline{V}^\dagger_\ell)_{\ell j} = \hat{y^\ell}_i \delta_{ij}$. In which the structure of the Yukawa coupling given in (8) is well reproduced by imposing constraints on the unitary matrices, $|(\overline{V}_{e,\ell})_{ij}| \lesssim \epsilon^{|X^{e,\ell}_i - X^{e,\ell}_j|}$ for given diagonal Yukawa couplings, $\hat{y^\ell}_i \sim \mathcal{O}(\epsilon^{X^e_i + X^\ell_i})$. Similar discussion is applied for quarks. These constraints are directly translated into the probable magnitudes of mass-insertion parameters at the electro-weak scale defined like, $(\delta^\ell_{RL})_{ij} \equiv A^\ell_{ij} v_d / \sqrt{m^{2(\widetilde{e})}_{ii} m^{2(\widetilde{\ell})}_{jj}}$, $(\delta^d_{RR}) \equiv m^{2(\widetilde{d})}_{ij} / \sqrt{m^{2(\widetilde{d})}_{ii} m^{2(\widetilde{d})}_{jj}}$ and $(\delta^d_{LL}) \equiv m^{2(\widetilde{q})}_{ij} / \sqrt{m^{2(\widetilde{q})}_{ii} m^{2(\widetilde{q})}_{jj}}$ in the SCKM basis [5]. Note that if some of the effective charges are degenerate, unitarity relation like $(\overline{V}_e)_{ik} X^e_k (\overline{V}^\dagger_e)_{kj} = (\overline{V}_e)_{i1} (\overline{V}^*_e)_{j1} (X^e_1 - X^e_2) - (\overline{V}_e)_{i3} (\overline{V}^*_e)_{j3} (X^e_2 - X^e_3) + \delta_{ij} X^e_2$ can dramatically suppresses the mixing from the naive estimation. This suppression mechanism is quite natural if underlying physics quantizes the kink masses in some unit, which originate from $U(1)_{FI}$ or graviphoton charges in the supergravity formulation [6].

Table 1: Lepton mass hierarchy vs constraint from $\mu \to e\gamma$. $(\Delta^\ell_{RL(LR)})_{12} \equiv (\delta^\ell_{RL(LR)})_{12}/4.8 \times 10^{-6}$.

$\hat{y}^\ell_i/\hat{y}^\ell_3 = \mathcal{O}(\epsilon^5, \epsilon, 1)$				$\hat{y}^\ell_i/\hat{y}^\ell_3 = \mathcal{O}(\epsilon^6, \epsilon, 1)$			
$X^e_i - X^e_3$	$X^\ell_i - X^\ell_3$	$\lvert(\Delta^\ell_{RL})_{12}\rvert$	$\lvert(\Delta^\ell_{LR})_{12}\rvert$	$X^e_i - X^e_3$	$X^\ell_i - X^\ell_3$	$\lvert(\Delta^\ell_{RL})_{12}\rvert$	$\lvert(\Delta^\ell_{LR})_{12}\rvert$
$(5,1,0)$	$(0,0,0)$	3.2	0.015	$(6,1,0)$	$(0,0,0)$	0.80	0.040
$(0,0,0)$	$(5,1,0)$	0.015	3.2	$(0,0,0)$	$(6,1,0)$	0.040	0.80
$\hat{y}^\ell_i/\hat{y}^\ell_3 = \mathcal{O}(\epsilon^5, \epsilon^2, 1)$				$\hat{y}^\ell_i/\hat{y}^\ell_3 = \mathcal{O}(\epsilon^6, \epsilon^2, 1)$			
$X^e_i - X^e_3$	$X^\ell_i - X^\ell_3$	$\lvert(\Delta^\ell_{RL})_{12}\rvert$	$\lvert(\Delta^\ell_{LR})_{12}\rvert$	$X^e_i - X^e_3$	$X^\ell_i - X^\ell_3$	$\lvert(\Delta^\ell_{RL})_{12}\rvert$	$\lvert(\Delta^\ell_{LR})_{12}\rvert$
No surviving models with mild tuning.				$(6,2,0)$	$(0,0,0)$	3.2	0.60
				$(2,2,0)$	$(4,0,0)$	1.6	3.2
				$(4,0,0)$	$(2,2,0)$	3.2	1.6
				$(0,0,0)$	$(6,2,0)$	0.60	3.2

We have explored various lepton flavor violating (LFV) and FCNC processes and find $\mu \to e, \gamma$ and ϵ_K give the most stringent constraint on the flavor structure of the SS SUSY breaking. Table 1 lists surviving set of $X^{e,\ell}_i$ from $\mu \to e, \gamma$ with $(\delta^\ell_{RL(LR)})_{12}$ divided by a value which saturates the current upper-bound, $BR(\mu \to e, \gamma) = 1.2 \times 10^{-11}$ for $|M_{1/2}| = 500$ GeV. The table shows that at least either X^e or X^ℓ should have degenerate charges to satisfy the bound. Table 2,3 show typical predictions of the models including other LFV processes. Note that degenerate charges of ℓ (e) yields the chirality of the processes opposite (similar) to the seesaw induced neutrino mass models [7]. On the other hand, table 4 shows the constraints from the CP violating parameter of $\overline{K}^0$-K^0 mixing, ϵ_K, where the mass-insertion parameters are divided by values saturating the observation $\epsilon_K = 2.282 \times 10^{-3}$ for $|M_{1/2}| = 500$ GeV. Models in the table can reproduce the quark masses and CKM angles with moderate tuning of the $\mathcal{O}(1)$ parameters, however, the ϵ_K constraint requires further $\mathcal{O}(10^{1-2})$ fine-tuning. However, if we allow abnormally

Table 2: Predictions of LFV rates for $y_i^\ell = \mathcal{O}(\epsilon^8, \epsilon^3, \epsilon^2)$ and $|M_{1/2}| = 500$GeV. $t_\beta \equiv \tan\beta$.

N_i^e	N_i^ℓ	$BR(\mu_R^+ \to e_L^+, \gamma)$	$BR(\tau_R^+ \to e_L^+, \gamma)$	$BR(\tau_R^+ \to \mu_L^+, \gamma)$
$6, 1, -2$	$2, 2, 2$	$9.2(1 + 0.11 t_\beta)^2 \times 10^{-12}$	$-$	$2.4(1 + 0.091 t_\beta)^2 \times 10^{-8}$
$8, 3, 2$	$-1, -1, -1$	7.8×10^{-12}	$-$	$2.4(1 + 0.091 t_\beta)^2 \times 10^{-8}$
N_i^e	N_i^ℓ	$BR(\mu_L^+ \to e_R^+, \gamma)$	$BR(\tau_L^+ \to e_R^+, \gamma)$	$BR(\tau_L^+ \to \mu_R^+, \gamma)$
$-1, -1, -1$	$8, 3, 2$	7.7×10^{-12}	$-$	$2.1(1 + 0.060 t_\beta)^2 \times 10^{-8}$
$2, 2, 2$	$6, 1, -2$	$7.7(1 + 0.075 t_\beta)^2 \times 10^{-12}$	$-$	$2.1(1 + 0.060 t_\beta)^2 \times 10^{-8}$

Table 3: Same as table 2 for $y_i^\ell = \mathcal{O}(\epsilon^8, \epsilon^4, \epsilon^2)$.

N_i^e	N_i^ℓ	$BR(\mu_R^+ \to e_L^+, \gamma)$	$BR(\tau_R^+ \to e_L^+, \gamma)$	$BR(\tau_R^+ \to \mu_L^+, \gamma)$
$2, 2, -2$	$6, 2, 2$	3.1×10^{-11}	3.4×10^{-9}	3.3×10^{-9}
$4, 4, 2$	$4, -2, -2$	$3.4(1 + 0.053 t_\beta)^2 \times 10^{-11}$	$3.7(1 + 0.047 t_\beta)^2 \times 10^{-9}$	$3.6(1 + 0.047 t_\beta)^2 \times 10^{-9}$
N_i^e	N_i^ℓ	$BR(\mu_L^+ \to e_R^+, \gamma)$	$BR(\tau_L^+ \to e_R^+, \gamma)$	$BR(\tau_L^+ \to \mu_R^+, \gamma)$
$4, -2, -2$	$4, 4, 2$	$3.1(1 + 0.032 t_\beta)^2 \times 10^{-11}$	$3.4(1 + 0.030 t_\beta)^2 \times 10^{-9}$	$3.3(1 + 0.030 t_\beta)^2 \times 10^{-9}$
$6, 2, 2$	$2, 2, -2$	3.1×10^{-11}	3.4×10^{-9}	3.3×10^{-9}

Table 4: Quark mass hierarchy vs ϵ_K.

N_i^q	N_i^d	$\Im(\delta_{LL}^d)_{12}^2/(1.5 \times 10^{-2})^2$	$\Im(\delta_{RR}^d)_{12}^2/(1.5 \times 10^{-2})^2$	$\Im[(\delta_{RR}^d)_{12}(\delta_{LL}^d)_{12}]/(2.2 \times 10^{-4})^2$
$3, 2, -1$	$3, 2, 2$	1.5×10^{-2}	7.1×10^{-1}	466
$3, 2, -1$	$4, 2, 2$	1.5×10^{-2}	2.8×10^{-2}	93
$3, 2, -1$	$3, 3, 2$	1.5×10^{-2}	2.8×10^{-2}	93
$3, 2, -1$	$4, 3, 2$	1.5×10^{-2}	5.7×10^{-3}	42

large or small $\mathcal{O}(1)$ parameters by one order of $\epsilon \approx 0.2$, we can eliminate dangerous right-handed mixing by choosing degenerate N_i^d and safely satisfy the constraint. We have also examined ΔM_K, $\Delta M_{B_{d,s}}$, ϵ'/ϵ , $b \to s, \gamma$, and CP asymmetry in $B_d \to \phi K_S$, however, can not find any meaningful constraints or predictions beyond the standard model.

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2. References

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